

Is the η Meson a Goldstone Boson?

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Abstract

The decoupling of the η meson from the nucleon, as recently deduced from analyzing $\bar{p}p$ collisions and η photoproduction off the proton at threshold, is shown to provide an argument against the octet Goldstone boson nature of the η meson. This argument concerns the structure of the strong isoscalar axial vector current. A vanishing ηN coupling means a vanishing contribution of the η pole term to the hypercharge nucleon axial vector current. Therefore, no partial conservation of the latter can be achieved. This new situation invalidates the octet Goldberger–Treiman relation for the ηN coupling constant that is indicative of the octet Goldstone boson nature of the η meson. In such a case there is no longer a compelling reason for the standard belief that the η meson should be coupled to the hadronic vacuum through the hypercharge axial vector current. Rather, it will be coupled to it as any ordinary neutral pseudoscalar meson via the neutral axial vector current of the electroweak theory. As a result of the suggested universal structure of the weak and strong neutral axial vector currents the η meson satisfies the strange analog of the Goldberger–Treiman relation so that it acquires features of a strange Goldstone boson. Hence the ηNN vertex constant appears proportional to Δs ,

the fraction of nucleon spin carried by the strange quark sea, thus explaining its suppression naturally.

I. STANDARD VIEW OF π^0 AND η MESONS

In the limit of zero masses of the up (u) and down (d) quarks, quantum chromodynamics (QCD) is invariant under global chiral rotations of the quark field, $q \rightarrow \exp(-\frac{i}{2}\vec{\alpha} \cdot \vec{\tau}\gamma_5)q$, $\bar{q} \rightarrow \bar{q}\exp(-\frac{i}{2}\vec{\alpha} \cdot \vec{\tau}\gamma_5)$. When chiral symmetry is broken spontaneously Goldstone bosons arise that can be exploited as effective degrees of freedom.

The simplest example for a chiral field theory is the Gell-Mann–Lévy σ -model [1] with a Mexican hat-type potential for a (fictitious) scalar σ -meson and the pions. Its infinitely degenerate ground states lie in the rim of the hat and are marked by a corresponding non-zero vacuum expectation value $\langle \sigma \rangle_0$ of the σ -meson. Massless current quarks become massive (constituent) with $m_q \sim \langle \sigma \rangle_0$ in the broken symmetry phase when one of the vacua is spontaneously selected. The rim is flat in the pion directions which means that pions still have zero mass. They are the corresponding Goldstone bosons [2]. More realistic descriptions of chiral hadron systems are provided by Nambu–Jona-Lasinio models that are patterned after superconductivity [3]. The effective degrees of freedom are Goldstone bosons and dynamical quarks whose momentum dependent mass $m_q(p^2)$ originates from a gap equation. Upon approximating the dynamical mass by $m_q(0) \approx m_N/3$ one can introduce the concept of a constituent quark of the nonrelativistic quark model (NQM) at low momentum.

The Noether currents of the two-flavor chiral group $SU(2)_L \otimes SU(2)_R$ coincide with the isovector axial vector currents of the Glashow–Weinberg–Salam (GWS) electroweak gauge theory

$$J_{\mu,5}^{(i)} = \bar{q}\gamma_\mu\gamma_5\frac{\lambda^i}{2}q, i = 1, 2, 3. \quad (1)$$

They couple the corresponding Goldstone bosons, the pions, to the hadronic vacuum in a Lorentz invariant way, e.g. for the neutral pion

$$\langle 0 | \bar{q} \gamma_\mu \gamma_5 \frac{\lambda^3}{2} q | \pi^0 \rangle = f_{\pi^0} m_{\pi^0} i k_\mu, \quad (2)$$

where k_μ is the pion four-momentum, and $f_{\pi^0} m_{\pi^0} = 85 \pm 0.3$ MeV is the *ordinary* weak decay constant of the neutral pion [4], while f_{π^0} stands for the respective *dimensionless* decay constant. Such non-vanishing matrix elements are signals of spontaneous chiral symmetry breakdown. They also enter into the weak decays of pions, e.g. the $\pi^+ \rightarrow \mu^+ + \nu_\mu$ decay as well as the π^0 decay into two photons through the neutral axial vector current at one corner of an intermediate quark triangle diagram with photons at the other two vertices, a mechanism known as chiral anomaly [5]. It shows up in the divergence of the neutral axial vector current for massless quarks as

$$\partial^\mu J_{\mu,5}^{(3)} = \frac{n_3 \alpha}{2\pi} F_{\mu\nu} \tilde{F}^{\mu\nu}, \quad (3)$$

where $\tilde{F}^{\mu\nu} = \frac{1}{2} \epsilon^{\mu\nu\rho\sigma} F_{\rho\sigma}$ is the dual of the electromagnetic field tensor, and $n_3 = \text{Tr}(\tau^3 Q^2) = 3[(\frac{2}{3})^2 - (-\frac{1}{3})^2] = 1$ is the charge factor corresponding to u and d triangles including colors.

The generalization to $SU(3)_L \otimes SU(3)_R$ chiral rotations in three flavor quark space (u,d,s) as symmetry transformations of the QCD Lagrangian describing low mass hadron systems is one of the basic paradigms of contemporary hadron physics that has its roots partly in the success of the quark model with the SU(3) flavor octets. There are then three different neutral axial vector Noether currents corresponding to chiral rotations [6]. One additional Noether current is the hypercharge axial vector current, $J_{\mu,5}^{(8)}$, that contains λ^8 . It is expected to couple to the η meson similar to Eq. (2) thereby defining the dimensionless η decay coupling constant f_η .

In contrast, the Noether current associated with the $U_A(1)$ symmetry of the QCD Lagrangian is the singlet axial vector current with $\lambda^0 = \sqrt{\frac{2}{3}} \hat{1}$ in Eq. (1) instead of λ^i . Its divergence is given by

$$\partial^\mu J_{\mu,5}^{(0)} = \frac{N_f \alpha_s}{4\pi} F_{\mu\nu}^c \tilde{F}_c^{\mu\nu}, \quad (4)$$

which has the $U_A(1)$ anomaly containing the gluon field strength $F_{\mu\nu}^c$ and its dual $\tilde{F}_c^{\mu\nu}$. In spite of the spontaneous breakdown of the $U_L(1) \otimes U_R(1)$ symmetry, it was suggested by 't

Hooft that no corresponding Goldstone boson arises because of instanton effects. Thus, the properties of the η' meson differ substantially from those of Goldstone bosons, and it is not considered here further.

In addition, a purely strange axial vector current $J_{\mu,5}^{(s)}$ can be defined through the ideal mixing between the hypercharge and singlet axial vector currents as

$$J_{\mu,5}^{(s)} = -\frac{1}{\sqrt{2}} \left(\sqrt{\frac{2}{3}} J_{\mu,5}^{(8)} - \sqrt{\frac{1}{3}} J_{\mu,5}^{(0)} \right) = \frac{1}{2} \bar{s} \gamma_\mu \gamma_5 s. \quad (5)$$

The matrix elements of the, say, isovector axial current for the nucleon state is parametrized as [7]:

$$\langle N | \bar{q} \gamma_\mu \gamma_5 \frac{\lambda^3}{2} q | N \rangle = \bar{\mathcal{U}}_N [g_A(q^2) \gamma_\mu + q_\mu G_2(q^2)] \gamma_5 \frac{\tau^3}{2} \mathcal{U}_N, \quad g_A(0) = \Delta u - \Delta d, \quad (6)$$

where $g_A(q^2)$ denotes the weak isovector axial form factor of the nucleon. Upon replacing λ^3 , g_A , and G_2 in Eq. (6) by λ^i , $g_A^{(i)}$, and $G_2^{(i)}$ with $i = 8, 0$, respectively, one obtains similar equations for the hypercharge and singlet axial form factors involving the hypercharge and singlet axial coupling constants

$$g_A^{(8)}(0) = \frac{1}{\sqrt{3}} (\Delta u + \Delta d - 2\Delta s), \quad g_A^{(0)}(0) = \sqrt{\frac{2}{3}} (\Delta u + \Delta d + \Delta s). \quad (7)$$

Here Δu , Δd , and Δs denote the fraction of proton spin carried by the u , d , and s quarks, respectively. Just for completeness, we introduce the nucleon matrix element of the purely strange axial vector current as

$$\langle N | \bar{s} \gamma_\mu \gamma_5 s | N \rangle = G_1^s(q^2) \bar{\mathcal{U}}_N \gamma_\mu \gamma_5 \mathcal{U}_N, \quad G_1^s(0) = \Delta s, \quad (8)$$

where G_1^s denotes the strange axial coupling of the nucleon.

If chiral symmetry is spontaneously broken so that Eq. (2) and its hypercharge analog are valid, the nucleon matrix elements of the flavor preserving isovector and hypercharge axial vector currents are in turn dominated by pion- and η meson pole terms in the soft limit $k_\mu \rightarrow 0$. This means that the π and η meson are the respective isovector and octet Goldstone bosons whose currents are partially conserved

$$\partial^\mu J_{\mu,5}^{(3)} = (f_{\pi^0} m_{\pi^0}) m_{\pi^0}^2 \phi_\pi^0, \quad \partial^\mu J_{\mu,5}^{(8)} = (f_\eta m_\eta) m_\eta^2 \phi_\eta. \quad (9)$$

These PCAC relations are clearly consistent with Eq. (2) and its hypercharge analog. As a consequence, the pseudoscalar coupling constants of the pion and the η meson to the nucleon (in turn denoted by $g_{\pi NN}$ and $g_{\eta NN}$) satisfy Goldberger–Treiman type relations and are in the ratio

$$\frac{g_{\eta NN}}{g_{\pi^0 NN}} = \frac{g_A^{(8)} m_N}{(f_\eta m_\eta)} : \frac{g_A m_N}{(f_{\pi^0} m_{\pi^0})} \approx \frac{g_A^{(8)}}{g_A}. \quad (10)$$

Here, the *ordinary* pion decay constant ($f_{\pi^0} m_{\pi^0}$) is taken to be approximately equal to the one of the η meson, $f_\eta m_\eta$. From Eq. (10) it follows that the size of the meson-nucleon coupling constant is another criterion to help one decide whether or not a meson is to be considered as a Goldstone boson and what kind of Goldstone boson it is.

While for the pions experimental data are well known to have confirmed the Goldberger–Treiman relation $f_\pi m_\pi g_{\pi NN} = g_A m_N$ to an accuracy of about 6% [8], this is not the case for $g_{\eta NN}$. The SU(3) flavor symmetry and the NQM predict a relatively large value for

$$g_{\eta NN} = \frac{1}{\sqrt{3}} \frac{3\mathcal{F} - \mathcal{D}}{\mathcal{F} + \mathcal{D}} g_{\pi^0 NN} = \frac{\sqrt{3}}{5} g_{\pi^0 NN}, \quad (11)$$

if the standard axial vector \mathcal{F}/\mathcal{D} ratio 2/3 is adopted. The same value is obtained when the NQM spin fractions $\Delta u = 4/3$, $\Delta d = -1/3$, $\Delta s = 0$ are used in Eqs. (6),(7), (10). On the other hand, the recent deep inelastic scattering experiment E143 [9] reports the following spin fractions for the proton

$$\Delta u = 0.84 \pm 0.05, \quad \Delta d = -0.43 \pm 0.05, \quad \Delta s = -0.08 \pm 0.05, \quad (12)$$

its analysis relying on the \mathcal{F}/\mathcal{D} axial vector ratio of the SU(3) flavor symmetry. If these spin fractions are used in Eqs. (6), (7), (10) in conjunction with $g_A = \mathcal{F} + \mathcal{D} = 1.2573 \pm 0.0028$, [4] then

$$g_{\eta NN} = (0.26 \pm 0.11) g_{\pi^0 NN} \quad (13)$$

results, again a fairly large value. With the spin fractions from [10]

$$\Delta u = 0.85 \pm 0.03, \quad \Delta d = -0.41 \pm 0.03, \quad \Delta s = -0.08 \pm 0.03, \quad (14)$$

one finds

$$g_{\eta NN} = (0.28 \pm 0.06)g_{\pi^0 NN}, \quad (15)$$

confirming the previous value with smaller error. Clearly, the more negative Δs is the larger $g_{\eta NN}$ will be if the η meson is taken to be the octet (or hypercharge) Goldstone boson.

The difference between these estimates for the octet coupling $g_{\eta NN}$ is about 20 to 25%. This is larger than the pion deviation quoted earlier but of the same order as the deviation, $\Delta_K = 1 - (m_N + m_\Lambda)g_A^{(4 \pm i5)}/(2f_K m_K g_{KN\Lambda}) \approx 30 \pm 15\%$, from the Goldberger-Treiman relation for the kaons [11] and in line with the mass scale $m_K \sim 3.5m_\pi$ to $m_\eta \sim 3.9m_\pi$. Using the NQM octet values for $g_A^{(8)}$ and $g_{\eta NN}$ in the corresponding deviation gives about 25% for the η meson. Corrections from chiral perturbation theory are of order 30% [12] and therefore much too small to help one understand better the problem of the suppressed ηNN coupling.

The size of the pseudoscalar η -nucleon coupling constant $g_{\eta NN}$ has been repeatedly subjected to comparisons with data. The analysis of $\bar{p}p$ collisions by means of a dispersion relation technique in [13] led to the surprising result of a vanishing η -nucleon coupling. Later on, fitting nucleon-nucleon as well as hyperon-nucleon phase shifts by means of the full Bonn potential, no need for an η meson exchange was found in [14], and [15]. Most recently, accurate differential cross sections for η photoproduction off protons at threshold obtained at the Mainz Microtron (MAMI) were shown to require an η -nucleon coupling constant suppressed by a factor of more than two relative to the quark model value in Eq. (11) giving ¹

¹ The small value for $g_{\eta NN}$ ensures cancellation between the Born terms and the background processes proceeding via ω and ρ intermediate states and leads to an almost 100% dominance of the S_{11} resonance in case of S-wave η photoproduction. This allows one to determine the N(1535)

$$|g_{\eta NN}| = 0.16g_{\pi NN}. \quad (16)$$

The violation of the octet Goldberger-Treiman relation is actually worse than a naive comparison of Eqs. (15) and (16) would suggest because Eq. (15) has to be compared to the tree level (contact) piece of $g_{\eta NN}$ in Eq. (16). We shall come back to this point at the end of the next section.

The comparison of the η -nucleon coupling constant predicted from the octet Goldberger-Treiman relation with the data seems to rule out the η meson as the octet Goldstone boson. Thus, one has to search for a mechanism that supports the decoupling of the η meson from the nucleon. In the next section we demonstrate that the suppressed ηN coupling constant can be understood if the η decay into the hadronic vacuum proceeds through the total neutral axial vector current associated with the Z boson rather than through the hypercharge axial vector current of the SU(3)-flavor quark model. Note that this total isoscalar axial vector current can be represented as an ideal mixing between weak currents having the same structure as the hypercharge and the flavor singlet axial currents. Within such a scheme the η - (and in addition the $f_1 NN$) vertex constants will appear proportional to Δs , the fraction of nucleon spin carried by the strange quark sea, thus explaining their suppression naturally. A summary of the main results is provided in Sect. III.

II. DECOUPLING OF THE η MESON FROM THE OCTET AXIAL VECTOR CURRENT

As a consequence of the Goldstone boson nature of the η meson its weak decay is usually supposed to proceed via the hypercharge axial vector current. This means that one ascribes to the η meson the fundamental ability to filter out its octet component from the isoscalar electroweak axial vector current, much like an optically active material filters out a particular

parameters from the $p(\gamma, \eta)p$ reaction at threshold. For details concerning the $g_{\eta NN}$ extraction procedure from data, the interested reader is referred to the original literature [16].

circularly polarized component from linearly polarized light. We question this viewpoint and argue that the experimentally observed decoupling of the η meson from the nucleon is naturally understood if one assumes that the η meson is coupled to the hadronic vacuum as an ordinary neutral pseudoscalar meson directly through the Z boson. The neutral gauge current (denoted by J_μ^Z) of the electroweak theory corresponding to Z boson exchange [17] is given by

$$J_\mu^Z = -2 \sum_{i=1}^3 \bar{\Psi}_{iL} \gamma_\mu \frac{\tau_3^{weak}}{2} \Psi_{iL} + 2 \sin^2 \theta_W J_\mu, \quad (17)$$

$$\Psi_{1L} = \begin{pmatrix} u \\ d \end{pmatrix}_L, \quad \Psi_{2L} = \begin{pmatrix} c \\ s \end{pmatrix}_L, \quad \Psi_{3L} = \begin{pmatrix} t \\ b \end{pmatrix}_L. \quad (18)$$

Here Ψ_{iL} denotes the i th lefthanded quark generation, $\tau_3^{weak}/2$ stands for the weak isospin, θ_W is the Weinberg angle, and J_μ the electromagnetic current. From Eq. (18) one immediately sees that the first quark generation acts simultaneously as a doublet both for weak and strong isospin. This observation explains the success of the current algebra statement on the identity between the weak and strong isovector axial vector currents. In contrast, the second and the third quark generations decompose into singlets with respect to strong isospin and so do their respective axial vector currents. The neutral axial gauge current of the electroweak theory (denoted by $J_{\mu,5}^{GWS}$) that emerges from the first term on the rhs in Eq. (18) reads

$$J_{\mu,5}^{GWS} = -\frac{1}{2} \bar{u} \gamma_\mu \gamma_5 + \frac{1}{2} \bar{d} \gamma_\mu \gamma_5 d + \frac{1}{2} \bar{s} \gamma_\mu \gamma_5 s + \dots = -J_{\mu,5}^{(3)} + J_{\mu,5}^{(s)}, \quad (19)$$

where the contribution of heavier flavors was ignored for simplicity. Therefore, the weak isosinglet axial current is purely strange. In the following we will adopt the viewpoint that the hypercharge and the singlet axial currents are ideally mixed, as are those of the electroweak theory and work out the decay matrix element of the η meson (considered as the octet scalar state $|\eta\rangle = \frac{1}{\sqrt{6}}(\bar{u}u + \bar{d}d - 2\bar{s}s)$ for simplicity). Subsequently, we will make use of a scheme that was first exploited by Jaffe in [18] to calculate of the decay constant, f_M , of a vector meson (M). The main assumption is that the coupling of a component $(\bar{q}_i q_i)_M$

of a 1^{--} vector meson M to a current $\bar{q}_j \gamma_\mu q_j$ is diagonal in flavor

$$\langle 0 | \frac{1}{2} \bar{q}_i \gamma_\mu q_i | (\bar{q}_j q_j)_M \rangle = \kappa_j(1^{--}) \delta_{ij} m_M^2 \epsilon_\mu^M. \quad (20)$$

In the spirit of flavor symmetry it is furthermore convenient to suggest flavor independence of κ_j at least for the light flavors, i.e., $\kappa_u(1^{--}) = \kappa_d(1^{--}) \equiv \kappa(1^{--})$. The decay constants of the vector mesons as calculated within Jaffe's scheme can be shown to equal to those predicted by the NQM.

A convenient extension of Jaffe's coupling scheme to the case of pseudoscalar mesons is obtained by replacing one mass power by the corresponding meson four-momentum. In doing so, the pion decay constant is found with the help of the relations

$$\langle 0 | J_{\mu,5}^{GWS} | \pi^0 \rangle = f_{\pi^0} m_{\pi^0} i k_\mu, \quad f_{\pi^0} = -\sqrt{2} \kappa(0^{--}), \quad (21)$$

where $\kappa_u(0^{--}) = \kappa_d(0^{--}) \equiv \kappa(0^{--})$ was assumed, and $|\pi^0\rangle = \frac{1}{\sqrt{2}}(\bar{u}u - \bar{d}d)$ used. In a similar way the decay constant f_η of the η meson is calculated as

$$\langle 0 | J_{\mu,5}^{GWS} | \eta \rangle = f_\eta m_\eta i k_\mu, \quad f_\eta = -\frac{2}{\sqrt{6}} \kappa_s(0^{--}). \quad (22)$$

In case the flavor symmetry is violated, the values of $\kappa_s(0^{--})$ and $\kappa_{u/d}(0^{--})$ may differ appreciably. Indeed, from the empirical values of $f_{\pi^0} m_{\pi^0} = 85 \text{ MeV}$ and $f_\eta m_\eta = 94 \text{ MeV}$ [4] one extracts $\kappa(0^{--}) = -0.44$, and $\kappa_s(0^{--}) = -0.21$, respectively, which corresponds to the measured ratio of

$$\frac{f_\eta m_\eta}{f_{\pi^0} m_{\pi^0}} \approx 1.1. \quad (23)$$

Although the chiral perturbation theory calculation of Ref. [19] predicts the value of $(f_\eta m_\eta)/(f_{\pi^0} m_{\pi^0}) \approx 1.3$ in nice agreement with data, the analysis given above shows that the closeness of the empirical $f_{\pi^0} m_{\pi^0}$ and $f_\eta m_\eta$ values is not necessarily indicative of SU(3) flavor symmetry but may just be brought about by the proper size of the coupling of the strange quarkonium to the isosinglet axial vector current. In case the hypercharge and the

singlet currents are ideally mixed, the isosinglet axial vector current of the nucleon coincides with the weak isosinglet axial vector current [20]²

$$\langle N | J_{\mu,5}^{(s)} | N \rangle = \frac{G_1^s}{2} \bar{\mathcal{U}}_N \gamma_\mu \gamma_5 \mathcal{U}_N, \quad G_1^s = \Delta s. \quad (24)$$

The universal structure of the neutral axial vector current in both the weak and strong interactions allows one to express the coupling constant of the η meson (at the tree level) to any target by means of its weak decay constant as [21]

$$\frac{g_{\eta NN}}{m_N} = \frac{G_1^s}{f_\eta m_\eta}. \quad (25)$$

Note that this equation is the 'strange' analog of the Goldberger-Treiman relation. Such expressions are commonly used in all (vector) meson dominance models and imply the standard assumption of unsubtracted dispersion relations for the form factors dominated by the corresponding poles. Therefore, Goldberger–Treiman relations in fact reflect current universality. For this deeper reason the quark model had suggested that the η meson weak decay proceed by the octet axial vector current, thus keeping universality of the latter at both the strong and weak vertices. With Eq.(25) the ratio $g_{\eta NN}/g_{\pi^0 NN}$ is calculated as

$$\left| \frac{g_{\eta NN}}{g_{\pi^0 NN}} \right| \approx \left| \frac{\Delta s f_{\pi^0} m_{\pi^0}}{g_A f_\eta m_\eta} \right| \leq \left| \frac{\Delta s}{\Delta u - \Delta d} \right| \sim 0.06, \quad (26)$$

where the values $\Delta s = -0.08$ and $g_A = 1.2573$ [10] have been used. Thus a natural explanation is found for the suppression of the *contact* (tree level) η nucleon coupling.

In view of the suppression of the tree level η -nucleon coupling, loop corrections containing non-strange mesons like the $a_0(980)\pi N$ triangular contribution both to the ηNN and $f_1 NN$ vertices acquire importance. Such corrections have first been considered in Refs. [21], [22] and shown to be extremely useful to achieve agreement with data on η photo-production off protons at threshold.

²Note that in Ref. [20] the amplitude of the isoscalar axial vector current is defined as $\frac{G_1^s}{4}$ with $G_1^s = 2\Delta s$.

III. CONCLUSIONS

In the present study we reviewed the status of the η meson as the *octet* Goldstone boson of the spontaneously broken $SU(3)_L \otimes SU(3)_R$ chiral symmetry. We showed that recent precision experimental MAMI data on η photoproduction off protons at threshold fail to confirm the *octet* Goldberger–Treiman relation for the pseudoscalar η -nucleon coupling constant. The smallness of the experimentally observed $g_{\eta NN}$ compared to its quark model value and another one implied by the proton spin fractions via $g_A^{(8)}$ was shown to have a natural explanation in assuming the η meson to couple to a neutral axial vector current having the same structure as the isosinglet axial current of the electroweak theory. Within such a scenario the ηNN coupling constant follows from the ‘strange’ Goldberger–Treiman relation and, being proportional to Δs , the fraction of nucleon spin carried by the strange quark sea, the *contact*, or tree level, ηNN vertex appears suppressed relative to its quark model value by the amount of $\Delta s/g_A^{(8)}$.

The suppression of the nucleon matrix element of the hypercharge axial current was considered in a series of papers with $SU(3)$ breaking [23] on the one hand, and in the context of both chiral perturbation theory [12] and the large N_c limit of QCD [24] on the other hand. In particular, in [24] it was found that in the large N_c limit $3\mathcal{F} - \mathcal{D} = 0.28 \pm 0.09$. It is easy to verify that our idea of the decoupling of the η meson from the hypercharge axial vector current does not contradict that small result provided it is properly interpreted. Indeed, in our notation $3\mathcal{F} - \mathcal{D}$ equals $\sqrt{3}g_A^{(8)}$. In case the η meson couples to the purely strange axial vector current only, $g_A^{(s)} := \Delta s$ replaces $g_A^{(8)}$ leading to $\sqrt{3}|\Delta s| \approx 0.14$. This is the value that has to be compared with 0.28 ± 0.09 above. It is important to remark that, if the reason for the suppressed $g_{\eta NN}$ coupling is the preferred coupling of the η meson to the strange axial vector current, then relating the \mathcal{F}/\mathcal{D} ratio to the ηN system lacks any justification. For this, the β decay of the Ξ^- hyperon is relevant. The value [25] $3(G_A/G_V)_{\Xi^- \rightarrow \Lambda} = 3\mathcal{F} - \mathcal{D} = 0.75 \pm 0.15$ from data [4] for this decay turns out to be much larger than the one associated with the experimentally determined $g_{\eta NN}$ coupling, so that

from this angle as well the η meson could not be the octet Goldstone boson.

The difference between Eq. (25) and the experimental value of $g_{\eta NN}^2/4\pi = 0.4$ in Eq. (16) can be attributed to triangular $a_0\pi N$ vertex corrections along the line of Ref. [21]. To recapitulate, our idea is to parametrize on the composite hadron level the suppression of the ηNN vertex in terms of the nucleon matrix element of the purely strange axial current, $\langle N | J_{\mu,5}^{(s)} | N \rangle$, that is subsequently increased by effective triangular vertices containing non-strange mesons thus accounting for the effective strangeness of the surrounding meson cloud necessarily originating from OZI violation effects. This aspect appears to be more rigorous than those considered in Refs. [23], [12] and [24] where the concept of the SU(3) flavor symmetry is still kept despite substantial violation effects. It is worthwhile to mention that, if the OZI rule applied to the 0^{--} mesons (as it does for the 1^{--} mesons), then the η meson would be purely strange and its coupling to the nucleon would be natural via the 'strange' Goldberger-Treiman relation given in Eq. (25). It is not so obvious to realize that, although mass formulas deduced from the constituent quark model prefer a small $\eta - \eta'$ mixing, the corresponding sea quark currents could be ideally mixed and the 'strange' Goldberger-Treiman relation in Eq. (25) valid that gives the η meson features of a 'strange' Goldstone boson.

As a further consequence of the scheme considered, additional information on the coupling constants of neutral axial vector mesons like the $f_1(1285)$ meson can be obtained: the f_1 nucleon coupling arises exclusively because of the violation of the OZI rule within the axial vector meson nonet, and the f_1 mesons can not be considered as the chiral partners of the ω and ϕ mesons from the vector meson nonet. Thus, the 1^{--} and 1^{++} nonets can no longer be viewed as complete chiral partners.

We also wish to stress that the decoupling of the η meson from the nucleon, as revealed by various experimental studies during the last decade, questions the validity of $SU(3)_L \otimes SU(3)_R$ chiral rotations as symmetry transformations of the quark flavor triplet. Rather, two-flavor chiral rotations, $SU(2)_L \otimes SU(2)_R$, seem to be valid that act on distinct I-, V- and U- spin quark doublets and give rise to flavor changing currents. These strong currents have the same

structure (up to a Cabibbo rotation within the U spin doublet) as the weak ones. Within this scenario the strong and weak isosinglet axial vector currents have equal structures too and don't contain any longer u and d quarks. Therefore, the η meson will couple to strange quarks only. Thus little room is left for η meson exchange among up and down quarks. Hence isospin-independent spin-spin and tensor forces between quarks should be attributed to instanton effects as considered in [26] rather than to η meson exchange. Finally, another point of this comment is that probing the strange content of hadrons by strongly interacting η mesons might be as important as by weakly interacting Z bosons, possibly opening a field for new experimental activities.

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